Afterglow Light Curve Modulated by a Highly Magnetized Millisecond Pulsar

Heon-Young Chang, Chang-Hwan Lee, and Insu Yi

Korea Institute for Advanced Study
207-43 Cheongryangri-dong Dongdaemun-gu, Seoul 130-012, Korea
e-mail: hyc@kias.re.kr, chlee@kias.re.kr, iyi@kias.re.kr

Abstract. We investigate consequences of a continuously energy-injecting central engine of gamma-ray burst (GRB) afterglow emission, assuming that a highly magnetized pulsar is left beaming in the core of a GRB progenitor. Beaming and continuous energy-injection are natural consequences of the pulsar origin of GRB afterglows. Whereas previous studies have considered continuous energy-injection from a new-born pulsar to interpret the deviation of afterglow light curves of GRBs from those with the simple power law behavior, a beaming effect, which is one of the most important aspects of pulsar emissions, is ignored in earlier investigations. We explicitly include the beaming effect and consider a change of the beaming with time due to a dynamical evolution of a new-born pulsar. We show that the magnitude of the afterglow from this fireball indeed first decreases with time, subsequently rises, and declines again. One of the most peculiar optical afterglows light curve of GRB 970508 can be accounted for by continuous energy injection with beaming due to a highly magnetized new-born pulsar. We discuss implications on such observational evidence for a pulsar.

Key words. gamma rays:bursts – pulsar:general – stars:magnetic fields

1. Introduction

Gamma-ray bursts (GRBs) are widely accepted to be produced when fast-moving, relativistic shells ejected from a central source in a relatively short period collide with slowly moving, yet relativistic shells that were ejected at an earlier time (Rees & Mészáros 1994; Paczyński & Xu 1994; Kobayashi et al. 1997; Daigne & Mochkovitch 1998, 2000). In connection with the so-called internal shock model, the external shock model also prevails as a possible origin of the GRB afterglows. In the external shock model the relativistic material is assumed to be decelerated via interactions with the surrounding medium. A shock wave results in heating the ambient matter to relativistic temperatures, and emitting photons in longer wave lengths (Rees & Mészáros 1992; Mészáros & Rees 1993; Paczyński & Rhoads 1993; Sari et al. 1996; Vietri 1997). GRBs and their afterglows seem to result from the dissipation of bulk energy in the relativistic outflows, which are in the form of a narrow beam rather than a spherical shell.

Even though the origin of the observed GRBs are still unknown, from the observations of several GRB afterglows the evidence of beamed GRBs has been accumulated (Sari et al. 1999; Halpern et al. 1999; Rhoads 1999).

There are several works on models for the geometry of GRBs (e.g., Chang & Yi 2001 and references therein) and their environments (e.g., Scalo & Wheeler 2001 and references therein). Much of the current research on GRBs is aimed at determining the nature and the origin of the central engine (Duncan & Thompson 1992; Narayan et al. 1992; Woosley 1993; Katz 1994; Usov 1992, 1994a; Shaviv & Dar 1995; Mészáros & Rees 1997b; Yi & Blackman 1997; Blackman & Yi 1998; Paczyński 1998; MacFadyen & Woosley 1999; Portegies Zwart et al. 1999; Li 2000; Wheeler et al. 2000; Zhang & Fryer 2001). Although the simple cosmological fireball afterglow model is in a good agreement with the observed light curves of a power law decay (e.g., Wijers et al. 1997; Waxman 1997a,b), the afterglow of GRB970228 observed by HST, for instance, shows a deviation from the simple power law behavior (Fruchter et al. 1997). Several works have been done to further investigate more subtle effects that can change the afterglow characteristics (Katz & Piran 1997; Mészáros & Rees 1997a, 1999; Rhoads 1999; Berger et al. 2000; Dai & Lu 2000; Kumar & Panaitescu 2000a, b; Panaitescu & Kumar 2000; Dai & Lu 2001).

Most fireball models assume that the energy injection into the fireball occurs in a short period of time compared with the lifetime of the afterglows (Rees & Mészáros 1998;

Kumar & Piran 2000; Sari & Mészáros 2000). However, in some types of central engines, such as a fast rotating new-born pulsar with the strong magnetic field (magnetar), a significant energy input into the fireball may in principle continue for a significantly longer timescale, and accordingly the temporal decay of the afterglow will be slower. Hence, it is worthwhile to investigate a continuously fed fireball in more details as a probe of the central engine of GRBs. We here consider the central engine that emits both an initial impulsive energy input E_{imp} and a continuous power. In fact, recently there were such attempts to provide an explanation for the deviation of the afterglow light curve from the simple power law (Dai & Lu 1998a, b; Zhang & Mészáros 2001a, b). Even though previous studies have considered continuous injection from a highly magnetized millisecond pulsar to interpret the deviation of afterglow light curves of some GRBs from the power law behavior, one of the most important aspects of pulsar emissions is ignored, that is, beaming. Beaming and continuous powering are the clear consequences of the pulsar origin of the GRB afterglows. In this Letter, this feature is included and a change of the beaming due to a dynamical evolution of the new born pulsar is explicitly taken into account. We find that the deviations suggested by previous studies indeed occur, but that the shape of the light curve is significantly modified due to a beaming effect. If such observational evidence for beaming is found, the corresponding pulsar origin for GRBs would imply that activities resulting in too massive objects to be a neutron star might be ruled out as a central engine of GRBs. Therefore, the question of whether or not a signature of beaming can be found in the afterglows thus extremely important.

2. Afterglows Continuously Powered by Pulsar

The total luminosity emitted from a young millisecond pulsar (MSP) has two important terms: an electromagnetic (EM) radiation term and a gravitational wave (GW) radiation term. Given that the spin-down is mainly due to electromagnetic dipolar radiation and to gravitational wave radiation, the spin-down law is given by

$$-I\Omega\dot{\Omega} = \frac{B_p^2 R^6 \Omega^4 \sin^2\alpha}{6c^3} + \frac{32G\epsilon^2 I^2 \Omega^6}{5c^5}, \qquad (1)$$

where Ω and $\dot{\Omega}$ are the angular frequency and its time derivative, B_p is the dipolar field strength at the magnetic poles, α is the angle between spin axis and the magnetic field axis, I is the moment of inertia, R is the pulsar radius, and ϵ is the ellipticity of the compact star (Shapiro & Teukolsky 1983). When an MSP rotates with spin frequency larger than a certain critical frequency, Ω_{crit}, nonaxisymmetric secular instability drives the pulsar into nonaxisymmetric configuration with nonzero quadrupole moment. This is the case where the GW luminosity dominates. In this case, the spin-down of the MSP occurs on a timescale $\tau_{\mathrm{gw}} = I_*\Omega_*^2/2L_{\mathrm{gw}} \sim 3 \times 10^{-3}\epsilon^{-2}I_{45}\Omega_4^{-4}$ s, where $I_* = I_{45}10^{45}$ g cm^2 is the

MSP moment of inertia, $\Omega_* = \Omega_4 10^4$ s^{-1} is the MSP spin frequency, and we have made use of the gravitational wave energy-loss rate, $L_{\mathrm{gw}} = 32G\epsilon^2 I_*^2\Omega_*^6/5c^5$. When an MSP rotates with spin frequency smaller than the certain critical frequency, the spin-down due to the electromagnetic dipole radiation gives a timescale $\tau_{\mathrm{em}} \sim I_*\Omega_*^2/2L_{\mathrm{em}} \sim 3 \times 10^2 I_{45}B_{15}^{-2}R_6^{-6}\Omega_4^{-2}$ s, where $R_* = R_6 10^6$ cm is the MSP radius, $B_* = B_{15}10^{15}$ G is the MSP dipole magnetic field, and we have assumed the electromagnetic dipole energy-loss rate, $L_{\mathrm{em}} = 2\mu_*^2\Omega_*^4/3c^3$, where $\mu_* = B_*R_*^3$ is the electromagnetic dipole moment. The decay solution $\Omega(T)$ as a function of time T includes both EM and GW emissions, but the energy input into the fireball is due to the EM dipolar emission L_{em} alone. Provided that the spin-down at a time is dominated by one or the other loss term, one can obtain approximate solutions. When GW radiation losses dominate the spin-down, we have $\Omega = \Omega_{0,\mathrm{gw}}(1 + T/\tau_{\mathrm{gw}})^{-1/4}$, or approximately $\Omega = \Omega_{0,\mathrm{gw}}$ for $T \ll \tau_{\mathrm{gw}}$, and $\Omega = \Omega_{0,\mathrm{gw}}(T/\tau_{\mathrm{gw}})^{-1/4}$ for $T \gg \tau_{\mathrm{gw}}$, where $\Omega_{0,\mathrm{gw}}$ being the initial spin frequency. When EM dipolar radiation losses dominate the spin-down, we have $\Omega = \Omega_{0,\mathrm{em}}(1 + T/\tau_{\mathrm{em}})^{-1/2}$, or approximately $\Omega = \Omega_{0,\mathrm{em}}$ for $T \ll \tau_{\mathrm{em}}$, and $\Omega = \Omega_{0,\mathrm{em}}(T/\tau_{\mathrm{em}})^{-1/2}$ for $T \gg \tau_{\mathrm{em}}$, where $\Omega_{0,\mathrm{em}}$ being the initial spin frequency. Assuming $\Omega_{\mathrm{crit}} > \Omega_{\mathrm{dynamo}}$, where Ω_{dynamo} is the spin required for dynamo action, there are two types of MSPs: a supercritical strong field rotator (SPS) with $\Omega > \Omega_{\mathrm{crit}} > \Omega_{\mathrm{dynamo}}$, and a subcritical strong field rotator (SBS) with $\Omega_{\mathrm{crit}} > \Omega > \Omega_{\mathrm{dynamo}}$. These two classes may give the bimodal distribution of short and long pulsar-induced GRBs (Yi & Blackman 1998). For an SPS pulsar, GW spin-down is important in initial period and promptly below Ω_{crit}. As soon as $\Omega < \Omega_{\mathrm{crit}}$ the spin-down becomes dominated by the EM losses. On the other hand, for an SBS pulsar, the spin-down is always dominated by EM regime.

The differential energy conservation relation for the self-similar blast wave can be written as $dE/dt = \mathcal{L}_0(t/t_0)^{q'} - \kappa'(E/t)$, where E and t are the energy and time measured in the fixed frame and q' and κ' are constants (Cohen & Piran 1999). The first term denotes the continuous luminosity injection, and the second term takes into account radiative energy losses in the blast wave. For $t > t_0$, the bulk Lorentz factor of the fireball scales with time as $\Gamma^2 \propto t^{-m}$, with m and κ' related by $\kappa' = m - 3$ (Cohen et al. 1998). If $m = 3$, it corresponds to the adiabatic case (Blandford & McKee 1976). In the observer frame, the time T is related to the fixed frame t by $dT = (1 - \beta)dt \simeq dt/2\Gamma^2$, and $T = \int_0^t (2\Gamma^2)^{-1}dt \simeq t/[2(m + 1)\Gamma^2]$ when $t \gg t_0$. The differential energy conservation relation in the observer frame is now given by $dE/dT = L_0(T/T_0)^q - \kappa(E/T)$, and can be integrated as

$$E = \frac{L_0}{\kappa + q + 1}\left(\frac{T}{T_0}\right)^q T + E_{\mathrm{imp}}\left(\frac{T}{T_0}\right)^{-\kappa}, T > T_0, \qquad (2)$$

where $T_0 = t_0/[2(m' + 1)\Gamma^2]$, m' being the self-similar index for $t < t_0$, $L_0 = 2\Gamma^2\mathcal{L}_0$, $q = (q' - m)/(m + 1)$, and

$\kappa = \kappa'/(m+1)$. Setting $T = T_0$, the total energy at the beginning of the self-similar expansion is the sum of two terms, $E_0 = L_0T_0/(\kappa+q+1)+E_{imp}$. The first term is the accumulated energy from the continuous injection before the self-similar solution begins. The second term E_{imp} is the energy injected impulsively by the initial blast.

The total energy of the blast wave given by Eq. 2 may be dominated either by the continuous injection term ($\propto T^{q+1}$) or by the initial impulsive term ($\propto T^{-\kappa}$), subject both to the relative values of the two indices and to the values of L_0 and E_{imp} (Dai & Lu 1998a, b; Zhang & Mészáros 2001a). One may classify three regimes according to the relative values of the two indices as discussed by Zhang & Mészáros (2001a). We are interested in the case where $q > -1-\kappa$, since otherwise a pulsar signature is no longer observable even if there is a pulsar in the central engine. If $q > -1-\kappa$, the first term in Eq. 2 will eventually dominate over the second term after a critical T_c. The injection-dominated regime begins at a critical time T_c defined by equating the injection and energy-loss terms in Eq. 2,

$$T_c = \text{Max}\left\{1, \left[(\kappa+q+1)\frac{E_{imp}}{L_0T_0}\right]^{1/(\kappa+q+1)}\right\}. \qquad (3)$$

If initially the impulsive term dominates, the critical time T_c could be much longer than T_0, depending on the ratio of E_{imp} and L_0T_0. The case where $T_c \geq T_0$ ensures that a self-similar solution has already formed when the continuous injection term dominates. Besides the condition of $q > -1-\kappa$, for detecting the signature of a pulsar the critical time T_c should be greater than τ_{em} or τ_{gw} as discussed below.

To obtain the temporal decay index of the afterglow light curve for which the MSP is responsible we adopt the cylindrical geometry instead of the spherical geometry, which may accommodate elongated beaming configurations. The rotational axis of the MSP coincides with the z-axis of the geometry. For a fireball blastwave decelerated by a homogeneous external medium with particle number density n, the energy conservation equation at time $t = r/c$ is given by

$$E = 2\pi R^2 rnm_pc^2\Gamma^2 = 2\pi R^2 ctnm_pc^2\Gamma^2, \ t > t_0, \qquad (4)$$

where r is the distance of the blastwave in z direction, R is a radius of the beaming which can be characterized by that of the light cylinder, c/Ω, Ω being given by the MSP evolution, and all other symbols have their usual meanings. One can work out the relationship between the temporal index α and the spectral index β, where $F_\nu \propto T^\alpha\nu^\beta$. To directly compare indices suggested by Zhang & Mészáros (2001a), we derived the scaling laws for the slow cooling regime assuming that the reverse shock is mildly relativistic (Sari et al. 1998). We derive the scaling law for $T \ll \tau_{em}$ as an illustrating example. The energy E in Eq. 2 should have the same time dependence as the first term on the right-hand side of Eq. 1, giving $T^{q+1} \propto t\Gamma^2$, or $m = -q/(q+2)$. Since in general,

Table 1. Temporal index α for the forward and the reverse shocks in various timescales.

	α^f	α^r
$T \ll \tau_{gw}$	$3+2q-q\beta^f$	$(9q+12-2q\beta^r)/4$
$T \gg \tau_{gw}$	$2+2q-\beta^f(q-1/2)$	$(18q+15-2(2q-1)\beta^r)/8$
$T \ll \tau_{em}$	$3+2q-q\beta^f$	$(9q+12-2q\beta^r)/4$
$T \gg \tau_{em}$	$1+2q-\beta^f(q-1)$	$(9q+3-2\beta^r(q-1))/4$

$\Gamma \propto t^{-m/2} \propto r^{-m/2} \propto T^{-m/2(m+1)}$ and $r \propto T^{1/(m+1)}$, the Γ dependence on T is given by $\Gamma \propto T^{q/4}$. The relationship between the temporal index α and the spectral index β can be obtained as below. For the forward shock, the synchrotron peak frequency $\nu_m^f \propto \Gamma^4 \propto T^{-2m/(m+1)} \propto T^q$, the peak flux $F_{\nu_m}^f \propto T^3\Gamma^8 \propto T^{(3-m)/(m+1)} \propto T^{3+2q}$. Thus, $\alpha^f = 3+2q-q\beta = (2m\beta^f+3-m)/(1+m)$ since $T^{\alpha^f}\nu_m^{\beta^f} \propto T^{3+2q} \propto T^{(3-m)/(m+1)}$. For the reverse shock, $\nu_m^r = \nu_m^f/\Gamma^2 \propto \Gamma^2 \propto T^{-m/(m+1)} \propto T^{q/2}$, the peak flux $F_{\nu_m}^r = \Gamma F_{\nu_m}^f \propto T^3\Gamma^9 \propto T^{(6-3m)/2(m+1)} \propto T^{(12+9q)/4}$. Thus, $\alpha^r = (2m\beta^r+6-3m)/2(1+m) = (12+9q-2q\beta)/4$. Similarly, one can derive the temporal index for each regime. In Table 1 we have summarized temporal indices in various regimes for the forward shock and the reverse shock, respectively.

If the continuous injection term becomes dominant over the impulsive term after T_c, the afterglow light curves rises after T_c and steepen again after some time, that is, about τ_{em} or τ_{gw}. Therefore, there may be two types of afterglow patterns for continuous injections according to a appropriate combination. Zhang & Mészáros (2001a) discussed conditions which allow to detect a signature for a pulsar and concluded that physical parameters are consistent with those of a magnetar in case of the afterglow features in case of GRB 000301c. In practice, however, τ_{gw} is very short, can be even shorter than T_c unless ambient matter density is very high, and is therefore unlikely to be observed. At around T_c and τ_{em}, q changes -1 to 0 and 0 to -2, respectively. These scaling laws represent a change from the standard adiabatic case to an EM-loss dominated regime as shown above (see also Zhang & Mészáros 2001a). For the forward shock, the temporal decay index changes around T_c from $3\beta^f/2$ to 3 and returns to $3\beta^f - 3$ after τ_{em}. For the reverse shock, the temporal decay index changes around T_c from $3\beta^r/4 - 3/8$ to 3 and returns to $3\beta^r/2 - 15/4$ after τ_{em}.

3. Discussions

A continuous energy injection signature in the GRB afterglow light curve may directly provide diagnostics about the nature of the injection as well as information on the GRB progenitor. Therefore, the question of whether or not a bump in the afterglow light curve is such an observational evidence for beaming from a pulsar is thus extremely important. In this sense, the correct geometry

should be applied to the model calculation. We have discussed the case where a pulsar is continuously injecting energy *cylindrically* rather than spherically. We suggest that a possible explanation for the deviation of the afterglow light curve of GRB 970508 is due to a beaming from a central pulsar. The optical afterglow of GRB 970508 has been explained as evidence of gravitational lensing event (Loeb & Perna 1998; Dado et al. 2001). So far the gravitational lensing can only account for the fact that the GRB afterglow shows a rise and a fall. We show in this Letter that with a correctly assumed geometry of energy injection the afterglow of GRB 970508 can be explained by a strongly magnetized fast-rotating pulsar with the continuous energy injection in a beam. In this calculation we ignore effects of the lateral expansion of the jet which may occur in the later time of the jet evolution. The light curve should be modified and becomes that of the spherical geometry case if the sideways expansion occurs in a timescale comparable to those we discussed after a neutron star formed. For instance, if the jet is expanded much faster at earlier stage than the light cone evolution, the cylindrical geometry effect becomes less obvious.

If the bump is caused by such a pulsar indeed, it puts constraints on the GRB progenitor models. That is, the corresponding pulsar origin for GRBs would imply that activities resulting in too massive objects to be a neutron star, such as, neutron star mergers, black hole formation, should be ruled out as a central engine of GRBs. During the birth of the neutron star, an initial fireball may occur through electromagnetic processes (Usov 1992). This has led to models in which GRBs are powered by rapidly spinning compact objects with strong magnetic fields (Blackman & Yi 1998; Blackman et al. 1996; Usov 1994b; Yi & Blackman 1997, 1998).

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